

Smooth Hybrid Inflation

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Abstract

We propose a variant of hybrid inflation which is applicable in a wide class of supersymmetric grand unified models. The observed temperature perturbations of cosmic background radiation can be reproduced with natural values of the parameters of the theory and a grand unification scale which is consistent with the unification of the minimal supersymmetric standard model gauge couplings as measured at LEP. The termination of inflation is smooth and does not produce any topological defects. Finally, we present a specific supersymmetric model where our smooth hybrid inflationary scenario is realized.

The usual realizations of the new and the chaotic inflationary scenarios⁽¹⁾ invoke a very weakly coupled gauge singlet scalar field known as inflaton. The extremely weak couplings of this field, which are necessitated by the smallness of the observed temperature fluctuations in the cosmic background radiation (CBR) are certainly unsatisfactory because they create an extra ‘hierarchy’ or ‘naturalness’ problem. Recently, Linde⁽²⁾ has proposed a new realization of the chaotic inflationary scenario based on a coupled system of two scalar fields one of which may not be a gauge singlet. The great advantage of this hybrid inflationary model is that it produces the observed temperature fluctuations in the CBR with natural values of the coupling constants. However, inflation terminates abruptly and is followed by a ‘waterfall’ regime during which topological defects can be easily produced. Consequently, in the cases where these defects include superheavy magnetic monopoles or domain walls we end up with an unacceptable cosmology. Of course, this is not the case with the original hybrid inflationary model⁽²⁾. Hybrid inflation has been adjusted⁽³⁾ so that it becomes applicable in a wide class of supersymmetric (SUSY) grand unified theories (GUTs) based on semi-simple gauge groups. The magnitude of the temperature fluctuations in CBR turns out to be directly related to the grand unification scale M_X . However, values of M_X consistent with the unification of the gauge couplings⁽⁴⁾ of the minimal supersymmetric standard model (MSSM), which is favored by recent LEP data, tend to give values of the temperature fluctuations in CBR considerably higher than the values allowed by observations. Also, the scheme heavily relies on radiative corrections.

In this paper we propose a variant of Linde’s potential which can also be

derived in a wide class of SUSY GUTs based on semi-simple gauge groups. Instead of the renormalizable trilinear coupling in the superpotential we utilize the first non-renormalizable contribution. All our analysis is done at the tree-level and no radiative corrections are needed. Although we get only a slight variation of Linde's potential, the cosmological scenario obtained is drastically different. Already since the beginning of inflation, the system follows a particular valley of minima which leads to a particular point of the vacuum manifold. So our inflationary scenario does not lead to production of topological defects. Also the termination of inflation is not as abrupt as in the hybrid case. It is quite smooth and resembles more the cases of new or chaotic inflation. The main advantage of our smooth inflationary scenario is that the measured value of the temperature fluctuations of CBR can be reproduced with natural values of the parameters and with a GUT scale, M_X , consistent with the unification of the MSSM gauge couplings. It is also remarkable that the scale controlling the non-renormalizable terms in the superpotential turns out to be of order 10^{18} GeV. The spectral index of density fluctuations is close to unity.

We consider a SUSY GUT based on a (semi-simple) gauge group G of rank ≥ 5 . G breaks spontaneously directly to the standard model (SM) group G_S at a scale $M_X \sim 10^{16}$ GeV. We assume that below M_X the only SM non-singlet states of the theory are the usual MSSM states in order for the successful MSSM predictions for $\sin^2\theta_w$ and α_s to be retained. The theory could also possess some global symmetries. The symmetry breaking of G to

G_S is obtained through a superpotential which includes the terms

$$W = s(-\mu^2 + \frac{(\phi\bar{\phi})^2}{M^2}). \quad (1)$$

Here $\phi, \bar{\phi}$ is a conjugate pair of left-handed SM singlet superfields which belong to non-trivial representations of the gauge group G and reduce its rank by their vacuum expectation values (vevs), s is a gauge singlet left-handed superfield, μ is a superheavy mass scale related to M_X , whereas M is a mass scale of the order of the ‘compactification’ scale $M_c \sim 10^{18}$ GeV which controls the non-renormalizable terms in the superpotential of the theory. The superpotential terms in eq.(1) are the dominant couplings involving the superfields $s, \phi, \bar{\phi}$ consistent with a continuous R-symmetry under which $W \rightarrow e^{i\theta}W, s \rightarrow e^{i\theta}s, \phi\bar{\phi} \rightarrow \phi\bar{\phi}$ and a discrete symmetry under which $\phi\bar{\phi}$ changes sign. Moreover, we assume that the presence of other SM singlets in the theory does not affect the superpotential in eq. (1) which is responsible for the non-zero vevs of the fields $\phi, \bar{\phi}$. This, if not automatic, is very easily achieved in the case of semi-simple gauge groups. The potential obtained from W in eq.(1), in the supersymmetric limit, is

$$V = |\mu^2 - \frac{(\phi\bar{\phi})^2}{M^2}|^2 + 4|s|^2 \frac{|\phi|^2|\bar{\phi}|^2}{M^4} (|\phi|^2 + |\bar{\phi}|^2) + D - terms, \quad (2)$$

where the scalar components of the superfields are denoted by the same symbols as the corresponding superfields. Vanishing of the D-terms is achieved with $|\bar{\phi}| = |\phi|$ (D-flatness condition). The supersymmetric vacuum

$$\langle s \rangle = 0, \quad \langle \phi \rangle \langle \bar{\phi} \rangle = \pm \mu M, \quad |\langle \bar{\phi} \rangle| = |\langle \phi \rangle| \quad (3)$$

lies on the D-flat direction $\bar{\phi}^* = \pm \phi$. Restricting ourselves to this particular direction and performing appropriate gauge, discrete and R-transformations

we can bring the complex $s, \phi, \bar{\phi}$ fields on the real axis, i.e, $s \equiv \frac{\sigma}{\sqrt{2}}$, $\bar{\phi} = \phi \equiv \frac{1}{2}\chi$, where σ and χ are real scalar fields. The potential in eq.(2) then takes the form

$$V(\chi, \sigma) = (\mu^2 - \frac{\chi^4}{16M^2})^2 + \frac{\chi^6 \sigma^2}{16M^4} \quad (4)$$

and the supersymmetric minima correspond to

$$| \langle \chi \rangle | = 2(\mu M)^{1/2}, \quad \langle \sigma \rangle = 0. \quad (5)$$

The mass acquired by the gauge bosons is $M_X = g(\mu M)^{1/2}$, where g is the GUT gauge coupling. The vacuum manifold of the theory, which is obtained from the minimum in eq.(5) by performing all possible gauge and global transformations, may have non-trivial homotopical properties in which case the theory predicts the existence of topological defects.

Let us now take a closer look at the potential in eq. (4). For any fixed value of σ , this potential, as a function of χ^2 , has a local maximum at $\chi^2 = 0$ and an absolute minimum lying at

$$\chi^2 \simeq \frac{4\mu^2 M^2}{3\sigma^2}, \quad \text{for } \sigma^2 \gg \mu M, \quad (6)$$

and $\chi^2 \simeq 4\mu M$, for $\sigma^2 \ll \mu M$. The value of the potential along the maxima at $\chi^2 = 0$ is constant, $V_{max}(\chi^2 = 0) = \mu^4$, whereas along the valley of minima, for $\sigma^2 \gg \mu M$, is $V_{min}(\sigma) \simeq \mu^4[1 - (2/27)(\mu^2 M^2/\sigma^4)]$.

We are now ready to turn to the discussion of the cosmological evolution of this system. We assume that after ‘compactification’ at a cosmic time $t_c \sim (\frac{3}{8\pi})^{1/2} \frac{M_P}{M_c^2}$, where $M_P = 1.2 \times 10^{19} \text{GeV}$ is the Planck mass, the universe emerges with energy density of order M_c^4 . We then, as usual, follow the development of a region of size $\ell_c \sim t_c$ where the scalar fields χ and σ happen

to be almost uniform with $|\sigma| \gg |\chi|$. Under these circumstances, one can easily see that initially the last term in eq.(4) is the dominant contribution to the potential energy of the system. Thus, the initial equations of motion for the χ and σ fields read

$$\ddot{\chi} + 3H\dot{\chi} + \frac{3\sigma^2\chi^5}{8M^4} \simeq 0 \quad (7)$$

and

$$\ddot{\sigma} + 3H\dot{\sigma} + \frac{\chi^6\sigma}{8M^4} \simeq 0, \quad (8)$$

with H being the Hubble parameter and dots denoting derivatives with respect to cosmic time. Assuming for the moment that σ remains approximately constant, we see that the frequency of oscillation of the χ field, which is of order $\chi_m^2|\sigma|/M^2$ (χ_m is the amplitude of the oscillations of χ), is much greater than the Hubble parameter $H \simeq (\pi/6)^{1/2}(\chi_m^3|\sigma|/M_P M^2)$, for $\chi_m \ll M_P$. So χ initially performs damped oscillations over the maximum at $\chi = 0$. For $|\sigma| \gg M_P$, the ‘frequency’ of oscillation of the σ field, which is of order χ_m^3/M^2 , is much smaller than H . Thus, our hypothesis that σ does not change much is justified. When the amplitude of the χ field drops to about $(\mu^2 M^2/|\sigma|)^{1/3}$, the μ^4 term dominates the potential in eq. (4) and the Hubble parameter becomes approximately constant and equal to $H = (8\pi/3)^{1/2}\mu^2/M_P$ and remains so thereafter till the end of inflation. However, the form of the eqs.(7) and (8) still holds till χ_m drops to about $\mu M/|\sigma|$. The ‘frequency’ of the σ field remains much smaller than H , for $(\mu^2 M^2/|\sigma|)^{1/3} \gtrsim \chi_m \gtrsim \mu M/|\sigma|$ ($|\sigma| \gg M_P$) and eq.(8) reduces to

$$3H\dot{\sigma} + \frac{\chi^6\sigma}{8M^4} \simeq 0. \quad (9)$$

As is easily shown, σ again remains almost constant within one expansion time, i.e., $|\Delta\sigma/\sigma| \ll \chi_m^6 M_P^2/\mu^4 M^4 \leq (M_P/|\sigma|)^2 \ll 1$. The frequency of the χ field $\sim \chi_m^2 |\sigma|/M^2$ remains greater than H , for $\chi_m \gtrsim \mu M/(M_P |\sigma|)^{1/2} \geq \mu M/|\sigma|$ and χ still performs damped oscillations. For smaller values of the χ field ($\mu M/(M_P |\sigma|)^{1/2} \gtrsim \chi \gtrsim \mu M/|\sigma|$), the Hubble parameter overtakes and χ enters into a ‘slow roll-over’ regime controlled by the equation

$$3H\dot{\chi} + \frac{3\sigma^2\chi^5}{8M^4} \simeq 0. \quad (10)$$

Eqs.(9) and (10) then imply that $3\sigma^2 - \chi^2 \simeq \text{constant}$. Thus, χ keeps dropping without any essential change in σ and, as one can deduce from eq.(10), reaches a value $\sim \mu M/|\sigma|$ in a cosmic time interval $\Delta t \sim (\sigma/M_P)^2 H^{-1} \gg H^{-1}$. After that, the χ^4 term in eq.(4) comes into play and the equation of motion of the χ field becomes

$$\delta\ddot{\chi} + 3H\delta\dot{\chi} + m_\chi^2(\sigma)\delta\chi \simeq 0, \quad (11)$$

where $\delta\chi$ is the deviation of χ from the minima in eq.(6) and $m_\chi^2(\sigma) \simeq 4\mu^4/3\sigma^2$ is the σ -dependent mass squared of the χ field along the valley of these minima. For $|\sigma| \gg M_P$, $m_\chi(\sigma) \ll H$ and eq. (11) reduces to

$$3H\delta\dot{\chi} + m_\chi^2(\sigma)\delta\chi \simeq 0. \quad (12)$$

Thus, χ reaches the minimum in a time interval

$$\Delta t \sim 6\pi\left(\frac{\sigma}{M_P}\right)^2 H^{-1} \quad (13)$$

within which the σ field stays essentially unaltered as one can see from eq.(9). To summarize, so far, we have seen that within a time interval given in eq.(13)

the χ field falls into the valley of minima in eq.(6) and relaxes at the bottom of this valley whereas the σ field still remains unchanged and much greater than M_P . Of course, this does not automatically mean that χ follows the valley of minima at subsequent times but, as we shall soon see, it turns out that it actually does.

One could also reach the same conclusion starting with field values considerably smaller than M_P . This requires that the initial energy density be $\sim \mu^4$ for σ to be slowly varying. Consider the situation where the initial value of σ is close to M . It is then easy to show that, for initial values of χ much smaller than $(4\sqrt{3\pi}\frac{(\mu M)^{1/2}}{M_P})^{1/3}(\mu M)^{1/2}$ but still greater than the value given in eq.(6), the χ field performs damped oscillations and ends up at the bottom of the valley of minima with σ remaining essentially unaltered.

To discuss the further evolution of the system let us, for the moment, suppose that χ , at subsequent times, follows the valley of minima in eq. (6). Eq. (9) then reads

$$3H\dot{\sigma} + \frac{8\mu^6 M^2}{27\sigma^5} \simeq 0, \quad (14)$$

which gives

$$\sigma \simeq \left(\frac{2\mu M M_P}{3\sqrt{2\pi}}\right)^{1/3} N^{1/6}(\sigma), \quad (15)$$

where $N(\sigma) \equiv H\Delta t$ is the number of e-foldings from the moment at which the σ field has the value σ till the end of inflation. Inflation ends when the ‘frequency’ of the σ field becomes smaller than $3H/2$, i.e., at

$$\sigma \simeq \sigma_o \equiv \left(\frac{2M_P}{9\sqrt{\pi}(\mu M)^{1/2}}\right)^{1/3}(\mu M)^{1/2} \gtrsim (\mu M)^{1/2}. \quad (16)$$

For $|\sigma| > M_P$, the evolution of χ is governed by eq.(12) and the time needed to reach its minimum (see eq.(13)) is much smaller than $N(\sigma)H^{-1}$ from

eq.(15) which is the time required for σ to change significantly. Also the quantum fluctuations of χ in de Sitter space, $\delta\chi \simeq H/2\pi$, are not strong enough to surpass the barrier at $\chi = 0$ if $\sigma \ll \sqrt{2\pi}(M/\mu)M_P$. For $\sigma \lesssim M_P$, $m_\chi \gtrsim H$, the quantum fluctuations of χ stop and we go back to eq.(11) which describes damped oscillations about the minimum in eq.(6) with damping time H^{-1} much smaller than $N(\sigma)H^{-1}$. Thus, we conclude that our starting assumption that χ follows the valley of minima is correct. Note that, due to supersymmetry breaking, the χ and σ fields acquire additional masses of order $M_s \sim 1$ TeV. It is easy to check that this mass contribution is never important for the evolution of the χ field. However, it plays a role for the σ field if its initial value happens to be greater than $\sigma_s \equiv (2/\sqrt{6})(M/M_s)^{1/3}\mu$. This introduces only unimportant complications to the above discussion which we will skip. In any case, as soon as σ drops below σ_s its SUSY breaking mass becomes subdominant.

The contribution of the scalar metric perturbation to the microwave background quadrupole anisotropy (scalar Sachs-Wolfe effect) is given⁽⁵⁾ by

$$\left(\frac{\Delta T}{T}\right)_S \simeq \left(\frac{32\pi}{45}\right)^{1/2} \frac{V^{3/2}}{M_P^3(\partial V/\partial\sigma)}|_{k\sim H} = 9\left(\frac{\pi}{10}\right)^{1/2} \frac{\sigma^5}{M_P^3 M^2}|_{k\sim H}, \quad (17)$$

where the right-hand side is evaluated at the value of the σ field where the length scale k^{-1} , which corresponds to the present horizon size, crossed outside the de Sitter horizon during inflation. The derivative $(\partial V/\partial\sigma)$ is, of course, calculated on the valley of minima of eq.(6) and substituting σ in terms of $N(\sigma)$ from eq.(15) and μ in terms of M_X , we obtain

$$\left(\frac{\Delta T}{T}\right)_S \simeq 5^{-1/2} 6^{1/3} \pi^{-1/3} N_H^{5/6} g^{-10/3} M_X^{10/3} M_P^{-4/3} M^{-2}, \quad (18)$$

where N_H is the number of e-foldings of the present horizon size during inflation. Taking $N_H = 60$, $M_X = 2 \times 10^{16}$ GeV, $g = 0.7$ (consistent with the MSSM unification) and $(\Delta T/T) \simeq 5 \times 10^{-6}$ from COBE, we get $M \simeq 9.4 \times 10^{17}$ GeV and $\mu \simeq 8.7 \times 10^{14}$ GeV. We ignored the gravitational wave contribution $(\frac{\Delta T}{T})_T$ to the quadrupole anisotropy since it turns out to be utterly negligible relative to the scalar component:

$$(\frac{\Delta T}{T})_T \simeq 0.78 \frac{V^{1/2}}{M_P^2} \simeq 4.1 \times 10^{-9}. \quad (19)$$

The amplitude of the density fluctuations on a given length scale k^{-1} as this scale crosses inside the postinflationary horizon is proportional to $N_{k^{-1}}^{5/6} \simeq N_H^{5/6} [k^{-1}(Mpc)/10^4]^{5/6 N_H}$, where $N_{k^{-1}}$ is the number of e-foldings of this scale during inflation (see eq.(18)). The spectral index is then given by

$$n = 1 - \frac{5}{3N_H} \simeq 0.97 \text{ (for } N_H \simeq 60), \quad (20)$$

which is very close to the Harrison-Zeldovich value ($n = 1$) and lies in the central range of values preferred by observations. We can also estimate the value of the σ field at the end of inflation (see eq.(16)) $\sigma_o \simeq 1.1 \times 10^{17}$ GeV $\simeq 5.4 M_X$ and its value when the present horizon size crossed outside the inflationary horizon (see eq.(15)) $\sigma_H \simeq (9N_H/2)^{1/6} \sigma_o \simeq 2.54 \sigma_o \simeq 2.7 \times 10^{17}$ GeV. The total number of e-foldings during inflation (see eq.(15)) turns out to be greater than 10^{12} , for initial $|\sigma| \geq M_P$, and SUSY breaking becomes irrelevant if the initial value of the σ field does not exceed $\sigma_s \simeq 5.8 M_P$ (for $M_s \simeq 1$ TeV).

The termination of inflation is not as abrupt as in the hybrid inflationary scenario where the slow roll-over regime is followed by the ‘waterfall’ . In

our case, after the end of inflation the σ and χ fields enter smoothly into an oscillatory phase about the global supersymmetric minimum of the potential in eq.(5) as in the ‘new’ inflationary scenario. The frequency of the oscillating fields is $m_\sigma = m_\chi = 2\sqrt{2}(\mu/M)^{1/2}\mu \simeq 7.6 \times 10^{13}$ GeV. These fields should eventually decay into lighter particles and ‘reheat’ the universe so that its baryon asymmetry can be subsequently produced. We will postpone until later the discussion of these important issues because it depends strongly on the details of the particle physics model one adopts.

Our inflationary scenario, in contrast to the hybrid one, produces no topological defects in the universe. This is because the σ and χ fields during inflation follow everywhere the same valley of minima which leads to a specific point of the vacuum manifold of the theory. As we discussed earlier, the χ field stays close to the bottom of the valley and is unable to fluctuate over the maximum at $\chi = 0$. Also, for $|\sigma| \lesssim M_P$, its mass is greater than H and, consequently, suffers no quantum fluctuations in de Sitter space. In hybrid inflation, the choice of the vacuum is made at the ‘waterfall’ after the end of inflation and, thus, topological defects can be easily produced.

It is well-known that, as soon as one replaces global by local supersymmetry, the potential of the theory changes drastically and inflation becomes, in general, impossible. This is, to a large extent, due to the generation of a mass for the inflaton which is larger than H . In our model, one can easily check that no such mass is generated provided we employ the canonical form of the *Kähler* potential. This fact allows us to hope that, although, our discussion is not expected to remain unaltered, the modifications necessitated by the inclusion of supergravity will not destroy the whole picture.

To construct an example of a SUSY model with the usual successful MSSM predictions of $\sin^2\theta_w$ and α_s where our smooth hybrid inflationary scenario is naturally realized, we consider a model with gauge symmetry group the subgroup $G \equiv SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times U(1)_T$ of E_6 and a global symmetry group $C \times B_1 \times B_2 \times R$ commuting with G . C is a Z_2 matter parity, B_1 and B_2 are Z_4 symmetries and R is a $U(1)$ R -symmetry under which the superpotential $W \rightarrow e^{i\theta}W$. The coupling constants of the five factors in G are assumed to be equal at a scale $M_c \sim 10^{18}\text{GeV}$, probably related to a more fundamental theory. Let $q, q^c, \ell, \ell^c, h, g, g^c, N$ denote chiral superfields transforming under G as $(3, 2, 1, +1/3, +1)$, $(\bar{3}, 1, 2, -1/3, +1)$, $(1, 2, 1, -1, +1)$, $(1, 1, 2, +1, +1)$, $(1, 2, 2, 0, -2)$, $(3, 1, 1, -2/3, -2)$, $(\bar{3}, 1, 1, +2/3, -2)$ and $(1, 1, 1, 0, +4)$ respectively. Also let s denote a gauge singlet. The superfield content of the model consists of three q_i , three q_i^c , three ℓ_i , three ℓ_i^c , one h , one \bar{h} , one g , one \bar{g} , one g^c , one \bar{g}^c , one ℓ_o^c , one $\bar{\ell}_o^c$, one ℓ_o , one $\bar{\ell}_o$, four $q_{(3+m)}$, four \bar{q}_m , four $q_{(3+m)}^c$, four \bar{q}_m^c , one N_o , one \bar{N}_o , three N_i , three \bar{N}_i and one s ($i = 1, 2, 3$ and $m = 1, \dots, 4$). Under C , all superfields remain invariant except $q_i, q_i^c, \ell_i, \ell_i^c$. Under the generator of B_1 , all superfields remain invariant except N_o which changes sign and $\ell_o^c, \bar{\ell}_o^c, g, g^c$ which are multiplied by i . Under the generator of B_2 , all superfields remain invariant except \bar{N}_o which changes sign and $\bar{h}, \bar{g}, \bar{g}^c$ which are multiplied by i . Finally under the $U(1)$ R -symmetry, all superfields have charge $+1/2$ except N_o, g, g^c which have charge $+1/3$, s which has charge $+1$ and $h, \ell_o^c, \bar{\ell}_o^c, \bar{N}_o$ which have charge 0 . With the above charge assignments, the $U(1)$ R -symmetry has color anomaly and therefore is a Peccei-Quinn symmetry solving the strong CP problem.

The gauge group $U(1)_T$ breaks down at a scale M_X through the vevs

of N_o, \bar{N}_o . The D -flat direction $|N_o| = |\bar{N}_o|$ is completely F -flat as well, because of the global symmetries. The large Yukawa couplings of the terms $gg^c N_o, \bar{h}^2 \bar{N}_o$ and $\bar{g}\bar{g}^c \bar{N}_o$ could lead to a radiative change of the sign of the SUSY breaking mass squared term for one linear combination of the N_o and \bar{N}_o bosons at a scale $\sim M_X$ thereby generating a large vev for this combination along the D -and F -flat direction. The gauge group $SU(2)_R \times U(1)_{B-L}$ breaks down to $U(1)_Y$ also at a scale M_X through the vevs of the neutral components ν_o^c and $\bar{\nu}_o^c$ in ℓ_o^c and $\bar{\ell}_o^c$. The relevant superpotential terms are the ones in eq.(1) with the role of ϕ and $\bar{\phi}$ played by ν_o^c and $\bar{\nu}_o^c$ respectively. The vev of N_o breaks B_1 down to a Z_2 subgroup, while \bar{N}_o breaks B_2 down to a Z_2 subgroup. Besides, the pair of $\langle N_o \rangle, \langle \bar{N}_o \rangle$ breaks the $U(1)_R$ -symmetry down to a Z_2 discrete symmetry. Starting with N_o and \bar{N}_o at their vevs we end up with the previously studied cosmological scenario.

We assume that the allowed tree-level mass terms for the gauge non-singlets are all $\sim M_X$. Below the scale M_X the only G_S -non-singlet states are the ones of the MSSM. Ordinary light quarks and leptons are contained in the $q_i, q_i^c, \ell_i, \ell_i^c$ fields all having negative matter parity. The electroweak Higgs doublets $h^{(1)}, h^{(2)}$ are the $SU(2)_R$ -doublet partners in h with mass $\sim \langle N_o^3 \bar{N}_o^2 \nu_o^c \bar{\nu}_o^c \rangle / M_c^6$. An immediate consequence of this simple structure is the relation $\tan\beta \simeq m_t/m_b$ among the vevs of $h^{(1)}, h^{(2)}$ and the 3rd generation quark masses. The g, g^c pair could, in principle, mediate proton decay but the global symmetries stabilize the proton almost completely.

Due to the appropriately chosen spectrum, the one-loop renormalization group equations above M_X predict identical running for the gauge couplings of all the factors in G . Combining this fact with their assumed equality at

M_c , we conclude that equality of the gauge couplings of G_S at the scale M_X is a justified boundary condition. The successful MSSM predictions for $\sin^2\theta_w$ and α_s follow immediately.

Returning to the cosmologically more interesting neutral sector we observe that we have exactly three right-handed neutrinos contained in the three ℓ_i^c 's which acquire Majorana masses $M_{\nu^c} \sim M^2\mu^2/M_c^3$ from the non-renormalizable terms $\nu_i^c\nu_j^c\nu_o^c(\bar{\nu}_o^c)^3/M_c^3$ ($i, j = 1, 2, 3$). Therefore the ν_i^c Majorana masses are at most $\sim 10^{12}$ GeV. The simple structure of neutrino Dirac mass terms m_ν^D leads to the simple relation $\tan\beta \simeq m_{\nu_\tau}^D/m_\tau$ for at least the 3rd generation. Thus, we expect a ν_τ mass $m_{\nu_\tau} \sim (\tan\beta m_\tau)^2/M_{\nu_\tau^c}$ which makes the ν_τ 's a significant component of the hot dark matter of the universe. The unbroken Z_2 matter parity C guarantees that the lowest supersymmetric particle remains stable and contributes to the cold component of the dark matter together with the axion.

The above mentioned terms providing the Majorana neutrino masses represent the dominant couplings leading to inflaton decay with width $\Gamma \sim \frac{1}{16\pi}M^3\mu^3M_c^{-6}m$, where $m \sim 10^{14}$ GeV is the inflaton mass. The corresponding 'reheat' temperature $T_r \sim \frac{1}{3}(\Gamma M_P)^{1/2} \sim 10^{11}$ GeV. The ν_i^c 's produced from the inflaton decay generate a lepton asymmetry⁽⁶⁾ which later will be transformed into the observed baryon asymmetry through the non-perturbative baryon-and-lepton-number-violating effects in the standard model at temperatures ~ 1 TeV.

The only topological defects predicted by this model are the domain walls associated with the spontaneous breaking of the discrete symmetries B_1 and B_2 down to their Z_2 subgroups. These symmetries are, at the first place,

broken by the vevs of N_o, \bar{N}_o . Our assumption that these fields acquire their vevs before inflation then leads to the complete absence of catastrophic domain walls in the universe.

In summary, in the context of SUSY GUTs based on semi-simple gauge groups, we constructed a smooth version of the hybrid inflationary scenario which can ‘naturally’ reproduce the observed temperature fluctuations in CBR with a GUT scale $\sim 10^{16}$ GeV consistent with the MSSM unification and a ‘compactification’ scale $\sim 10^{18}$ GeV. This inflationary scenario does not produce any topological defects. We also presented a specific model where our smooth hybrid inflation is realized.

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